

# A Novel Mechanism for $J/\psi$ Disintegration in Relativistic Heavy Ion Collisions

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In this paper we discuss the possibility of  $J/\psi$  disintegration due to  $Z(3)$  domain walls that are expected to form in a QGP medium. These domain walls give rise to localized color electric field which disintegrates  $J/\psi$ , on interaction, by changing its color composition and simultaneously exciting it to higher states of  $c\bar{c}$  system.

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## I. INTRODUCTION

The ongoing relativistic heavy ion collision experiments at RHIC (BNL) and LHC (CERN) have provided very valuable insights in understanding certain aspects of QCD. One such aspect is the existence of a new phase of matter known as quark-gluon plasma (QGP). QGP is essentially the deconfined phase of QCD, where free quarks and gluons exist in thermal equilibrium. Matsui and Satz [1] proposed that due to the presence of this medium, potential between  $q\bar{q}$  is Debye screened, resulting in the swelling of quarkonia. If the Debye screening length of the medium is less than the radius of quarkonia, then  $q\bar{q}$  may not form bound states. This is the conventional mechanism of quarkonia disintegration. Due to this melting, the yield of quarkonia will be *suppressed*. This was proposed as a signature of QGP and has been observed experimentally [1, 2]. However, there are other factors too that can lead to the suppression of  $J/\psi$  because of which it has not been possible to use  $J/\psi$  suppression as a clean signal for QGP.

In this paper, we propose a novel mechanism of quarkonia disintegration via QCD  $Z(3)$  domain walls. These walls appear as topological defects due to spontaneous breaking of  $Z(3)$  symmetry in QGP [3–5]. The thermal expectation value of Wilson loop (Polyakov loop) acts as the order parameter for confinement-deconfinement phase transition taking zero value in the confining phase (corresponding to infinite free energy of a test quark) and a non-zero value in the QGP phase (with finite free energy of a test quark). Polyakov loop transforms non-trivially under  $Z(3)$  transformations, hence its non-zero expectation value leads to spontaneous breaking of  $Z(3)$  symmetry in the QGP phase [6, 7]. With the possibility of realization of the QGP phase in RHIC and LHC experiments, we have the real opportunity to study topological domain walls, resulting from this spontaneous  $Z(3)$  symmetry breaking, in laboratory. The formation and evolution of these walls have been discussed in context of RHIC experiments [8, 9]. The associated  $QGP$  string

formation [10] has also been discussed by some of us. It is important to mention here that such topological defects invariably form during a phase transition. The formation process of topological defects is governed by formation of a sort of domain structure during a phase transition, and is usually known as the *Kibble mechanism* [11]. The network of defects formed depends on the details of phase transition only through the correlation length. In fact defect distribution, e.g. defect density, per correlation volume is universal and depends only on the symmetry breaking pattern and space dimensions.

Questions have been raised on the *reality* of these  $Z(3)$  domains. The existence of these  $Z(3)$  vacua becomes especially a non-trivial issue when considering the presence of dynamical quarks. The effect of quarks on  $Z(3)$  symmetry and  $Z(3)$  interfaces etc. has been discussed in detail in the literature [12, 13]. It has also been argued that the  $Z(3)$  symmetry becomes meaningless in the presence of quarks [12]. Other view-point, as advocated in many papers, asserts that one can take the effect of quarks in terms of explicit breaking of  $Z(3)$  symmetry [13–15]. We follow this approach and assume that the effects of dynamical quarks can be incorporated by introducing explicit symmetry breaking terms in the effective potential for the Polyakov loop. This makes  $Z(3)$  domain wall dynamical with pressure difference between the two different vacua being non-zero. This will lead to asymmetric profile of the Polyakov loop. In the present paper, we will ignore these *asymmetry* effects due to dynamical quarks, and will continue using  $Z(3)$  interfaces without any explicit symmetry breaking term. In a future work we will come back to include the effects of explicit symmetry breaking.

We mention recent lattice results in ref.([16]) which indicate strong possibility of the existence of these  $Z(3)$  domains at high temperatures in the presence of dynamical quarks. These results suggest that (metastable)  $Z(3)$  domains appear at temperatures above about 700 MeV. Though, we stress that it does not look appropriate to take these results as conclusive, especially the quantitative part. Thus one would like to consider the possibility that  $Z(3)$  vacua may persist for somewhat lower temperatures also as discussed in this paper. In any case, the mechanism discussed here provides additional source of disintegration for  $J/\psi$  even at high temperature. It is important to note that our mechanism will lead to disin-

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tegration of Upsilon also which will be relevant even at 700 MeV. We will present this study in a future work.

In case of early universe, these  $Z(3)$  walls can lead to baryon inhomogeneity generation [17]. It was shown in ref.[18] that background gauge field  $A_0$  associated with *generalized*  $Z(N)$  interfaces can lead to spontaneous CP violation in SM, MSSM and SUSY models, which, in turn, can lead to baryogenesis in the early universe. A detailed quantitative analysis of this spontaneous CP violation was done in [19], in the context of quark/antiquark scattering from  $Z(3)$  walls in the QGP phase. The main approach followed in refs.[18, 19] was based on the assumption that the profile of the Polyakov loop order parameter  $l(x)$  corresponds to a sort of condensate of the background gauge field  $A_0$  (in accordance with the definition of the Polyakov loop). This profile of the background gauge field can be calculated from the profile of  $l(x)$ . Such a gauge field configuration in the Dirac equation leads to different potentials for quarks and antiquarks, leading to spontaneous CP violation in the interaction of quarks and antiquarks from the  $Z(3)$  wall. This spontaneous CP violation was first discussed by Altes et al. [18, 20] in the context of the universe and in ref. [21] for the case of QCD. In [19], the profile of Polyakov loop  $l(x)$  between different  $Z(3)$  vacua was used (which was obtained by using specific effective potential for  $l(x)$  as discussed in [14, 15]) to obtain the profile of  $A_0$ . This background  $A_0$  configuration acts as a potential for quarks and antiquarks. It was shown in ref.[19] that the quarks have significantly different reflection coefficients than anti-quarks and the effect is stronger for heavier quarks. For a discussion of calculation of  $A_0$  profile, see ref. [19].

In this paper, we discuss the effect of this spontaneous CP violation on the propagation of quarkonia in the QGP medium, in particular, the  $J/\psi$  meson.  $J/\psi$  are produced in the initial stages of relativistic heavy ion collisions. As these are heavy mesons ( $m \sim 3\text{GeV}$ ), they are never in equilibrium with the QGP medium produced in present heavy-ion collision experiments. However, there are finite  $T$  effects (like Debye screening etc.) affecting its motion in a thermal bath. We ignore these effects initially and comment on it towards the end. Note that if the Debye length is larger, then the conventional mechanism of  $J/\psi$  melting does not work. As we will argue, for large Debye screening, our mechanism of  $J/\psi$  disintegration works better as any possible screening of the domain wall over the relevant length scale of  $J/\psi$  will be small. If a domain wall is present in the QGP, then a  $J/\psi$  moving through the wall will have a non-trivial interaction with it. Due to the CP violating effect of the interface on quark scattering,  $c$  and  $\bar{c}$  in  $J/\psi$  experience different color forces depending on the color of the quark and the color composition of the wall. This not only changes the color composition of  $c\bar{c}$  bound state (from color singlet to color octet state) but also facilitates its transition to higher excited states (for example  $\chi$  states). Color octet quarkonium states are unbound (also, the  $\chi$  state has

larger size than  $J/\psi$  and the Debye length), hence they will dissociate in the QGP medium. This summarizes the basic physics of our model discussed in this paper for quarkonia disintegration due to  $Z(3)$  walls.

The paper is organized in the following manner. In section II we discuss the interaction of  $J/\psi$  with the background gauge field  $A_0$  arising from the profile of  $l(x)$  and discuss its color excitations. Subsequently we consider spatial excitations of  $J/\psi$  and calculate the disintegration probability. Section III discusses results, and conclusions are presented in Sect.IV.

## II. INTERACTION OF $J/\psi$ WITH A $Z(3)$ WALL

In our model,  $J/\psi$  interacts with the gauge field  $A_0$  corresponding to the  $l(x)$  profile of the  $Z(3)$  wall. This allows for the possibility of color excitations of  $J/\psi$  as well as the spatial excitations of its wave function. First we discuss the possibility of color excitations of  $J/\psi$ . Subsequently, we will discuss spatial excitations of  $J/\psi$ .

### A. Color excitation of $J/\psi$

We work in the rest frame of  $J/\psi$  and consider the domain wall coming and hitting the  $J/\psi$  with a velocity  $v$  along  $z$ -axis. The gauge potential and coordinates are appropriately Lorentz transformed as

$$A_0(z) \rightarrow A'_0(z') = \gamma (A_0(z) - v A_3(z)) \quad (1a)$$

$$A_3(z) \rightarrow A'_3(z') = \gamma (A_3(z) - v A_0(z)) \quad (1b)$$

$$z = \gamma (z' + vt') \quad (1c)$$

We assume that there is no background vector potential,  $A_i(z) = 0$  ;  $i = 1, 2, 3$ .  $A'_3$  obtained from Eqn.. (1b) has only  $z'$  dependence, so it does not produce any color magnetic field. Further, using the non-relativistic approximation of the Dirac equation one can see that the perturbation terms in the Hamiltonian (say,  $H^1(A'_3)$ ) involving  $A'_3$  are suppressed compared to the perturbation term ( $H^1(A'_0)$ ) involving  $A'_0$  at least by a factor

$$\frac{H^1(A'_3)}{H^1(A'_0)} \sim \frac{v}{c} \frac{1}{m_c r_{J/\psi}} \quad (2)$$

where  $r_{J/\psi}$  is the size of the  $J/\psi$  wave function and  $m_c$  is the charm quark mass. As we will see, the largest value of  $v/c$  we consider is 0.20 - 0.24 (above which transition amplitude becomes too large to trust first order perturbation approximation). With  $r_{J/\psi} \simeq 0.4$  fm, the suppression factor in Eqn.(2) is of order 10 %. Thus we neglect perturbation due to  $A'_3$  and only consider perturbation due to  $A_0$  as given by Eqn.(1a). We use first order time dependent perturbation theory to study the excitation of  $J/\psi$  due to the background  $A_0$  profile and

consider the transition of  $J/\psi$  from initial energy eigenstate  $\psi_i$  with energy  $E_i$  to the final state  $\psi_j$  with energy  $E_j$ . The transition amplitude is given by

$$\mathcal{A}_{ij} = \delta_{ij} - i \int_{t_i}^{t_f} \langle \psi_j | \mathcal{H}_{int} | \psi_i \rangle e^{i(E_j - E_i)t} dt. \quad (3)$$

We take incoming quarkonia to be a color singlet state. The interaction of the quarkonia with the wall is written as

$$\mathcal{H}_{int} = V^q(z'_1) \otimes \mathbb{1}^{\bar{q}} + \mathbb{1}^q \otimes V^{\bar{q}}(z'_2) \quad (4a)$$

$$\text{with } V^{q,\bar{q}}(z'_{1,2}) = g A_0^{q,\bar{q}}(z'_{1,2}), \quad (4b)$$

where  $A_0^{q,\bar{q}}(z'_{1,2})$  is the background field configuration in the rest frame of  $J/\psi$ .  $z'_1$  and  $z'_2$  are the coordinates of  $q$  and  $\bar{q}$  in quarkonia and  $g$  is the gauge coupling. The gauge potential  $A_0$  is taken in the diagonal gauge as

$$A_0 = \frac{2\pi T}{g} (a\lambda_3 + b\lambda_8), \quad (5)$$

where  $\lambda_3$  and  $\lambda_8$  are the Gell-Mann matrices. Under CP,  $A_0 \rightarrow -A_0$ , hence  $A_0^{\bar{q}} = -A_0^q$ . Now, both the initial and the final states have a spatial, spin and color part. The incoming quarkonia is a color singlet while outgoing state could be a singlet or an octet. Using Eqn. (4), (5) and extracting only the color part of interaction, we get

$$\begin{aligned} \langle \psi_{out} | \mathcal{H}_{int} | \psi_{singlet} \rangle &= \langle \psi_{out} | g A_0^q(z'_1) \otimes \mathbb{1}^{\bar{q}} | \psi_{singlet} \rangle \\ &+ \langle \psi_{out} | \mathbb{1}^q \otimes g A_0^{\bar{q}}(z'_2) | \psi_{singlet} \rangle. \end{aligned} \quad (6)$$

The color singlet state of  $J/\psi$  is written as

$$\begin{aligned} |\psi_{singlet}\rangle &= \frac{1}{\sqrt{3}} \left[ \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}^q \otimes \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}^{\bar{q}} + \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix}^q \otimes \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix}^{\bar{q}} \right. \\ &\quad \left. + \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix}^q \otimes \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix}^{\bar{q}} \right]. \end{aligned} \quad (7)$$

If the outgoing state is also a singlet then, each term on RHS of Eqn. (6) is zero due to the traceless nature of  $A_0$ . Eqn. (3) gives  $\mathcal{A}_{ij} = 1$  for ground state ( $i = j$ ). (Meaning, one will then need to resort to 2nd order perturbation theory for consistency). For higher orbital states ( $i \neq j$ ), amplitude is identically zero. A color octet state like  $|r\bar{g}\rangle$ , can be written as

$$|r\bar{g}\rangle = \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}^q \otimes \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix}^{\bar{q}}. \quad (8)$$

For such an outgoing state each term on RHS of Eqn. (6) again vanishes identically because of the diagonal form of  $A_0$ , resulting in zero transition probability. Same argument leads to zero transition probability

to all other octet states with similar color content, viz.  $b\bar{g}$ ,  $b\bar{r}$ ,  $g\bar{r}$ ,  $\bar{g}b$ ,  $r\bar{b}$ . There are only two states with non-zero color contribution to transition probability. They are

$$|r\bar{r} - b\bar{b}\rangle = \frac{1}{\sqrt{2}} \left[ \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}^q \otimes \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}^{\bar{q}} - \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix}^q \otimes \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix}^{\bar{q}} \right] \quad (9)$$

and

$$\begin{aligned} |r\bar{r} + b\bar{b} - 2g\bar{g}\rangle &= \frac{1}{\sqrt{6}} \left[ \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}^q \otimes \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}^{\bar{q}} + \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix}^q \otimes \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix}^{\bar{q}} \right. \\ &\quad \left. - 2 \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix}^q \otimes \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix}^{\bar{q}} \right]. \end{aligned} \quad (10)$$

Using Eqn. (9) and (10) in conjunction with Eqn. (5), (1) and (6), we get the color part of transition probability as

$$\langle r\bar{r} - b\bar{b} | \mathcal{H}_{int} | \psi_{singlet} \rangle = \frac{1}{\sqrt{6}} (A_0^r - A_0^b) \quad \text{and} \quad (11a)$$

$$\langle r\bar{r} + b\bar{b} - 2g\bar{g} | \mathcal{H}_{int} | \psi_{singlet} \rangle = \frac{1}{\sqrt{18}} (A_0^r + A_0^b - 2A_0^g), \quad (11b)$$

where,  $A_0^r$ ,  $A_0^b$  and  $A_0^g$  are the diagonal components of the matrix  $A_0^q(z'_1) - A_0^{\bar{q}}(z'_2)$ . Eqn. (11a) and (11b) are the effective interactions that lead to the excitations of incoming  $J/\psi$  (in the color singlet state of  $c\bar{c}$ ) to the corresponding octet state. Due to repulsive Coulombic interaction of  $q$  and  $\bar{q}$  in the octet representation, one may expect that  $J/\psi$  may disintegrate while traversing through a  $Z(3)$  wall purely by color excitation. However, we will see in the next section that this is not so and one needs to also consider spatial excitation of  $J/\psi$  due to  $Z(3)$  wall.

## B. Spatial excitations of $J/\psi$

We now consider the spatial excitations. The spatial part of the states is decided by the potential between  $c\bar{c}$  in  $J/\psi$  which is taken as,

$$V(|\vec{r}_1 - \vec{r}_2|) = -\frac{\alpha_s C_F}{|\vec{r}_1 - \vec{r}_2|} + C_{cnf} \sigma |\vec{r}_1 - \vec{r}_2| \quad (12)$$

where  $\alpha_s$  is the strong coupling constant and  $\sigma$  is the string tension. For  $J/\psi$ , we will use charm quark mass  $m_c = 1.28$  GeV,  $\alpha_s = \pi/12$ , and  $\sigma = 0.16$  GeV<sup>2</sup> [22, 23].

$C_F$  is the color factor depending on the representation of the  $c\bar{c}$  state.  $C_F = 4/3$  for singlet state, while  $C_F = -1/6$  for the octet states showing the repulsive nature of the Coulombic part of the interaction for the octet states.  $C_{cnf}$  denotes the representation dependence of the confining part of the potential. For general sources, this factor follows Casimir scaling [24] for the string tension. For  $J/\psi$  in singlet representation,  $C_{cnf} = 1$  with the value of  $\sigma$  used here [22, 23]. It is not clear what should be the value of  $C_{cnf}$  if  $c\bar{c}$  are in the octet representation. As the Coulombic part of the potential is repulsive for the octet state of  $c\bar{c}$  (with  $C_F = -1/6$ ), it is not clear if there should be a confining part of the potential at all in this case for large distances. Early lattice simulations had indicated some possibility of mildly rising potential for the confining part for  $q\bar{q}$  in octet representation [25]. However, recent simulations do not show any such possibility. At large distances, the net potential between a  $q$  and  $\bar{q}$  in color octet state appears to be independent of distance [26]. With the repulsive Coulombic part, this implies a very small value for  $C_{cnf}$  for the confining part. For our purpose it suffices to assume that in the octet representation,  $J/\psi$  becomes unbound, having repulsive interaction at short distances.

We have seen above that the form of  $A_0$  in Eqn.(5) only allows for transition from color singlet to two of the color octet states given in Eqns.(9),(10). As we discussed above,  $c\bar{c}$  in color octet state is unbound. Thus our task should be to consider transition from initial color singlet  $J/\psi$  to unbound state of  $c\bar{c}$ , say in plane waves. However, this also does not look correct as the initial  $J/\psi$  (in the color singlet state) transforms to a color octet state only as it traverses the  $Z(3)$  wall (as coefficients  $a$  and  $b$  in Eqn.(5) undergo spatial variations). Thus during the early part of the passage of  $J/\psi$  through the wall, it should be dominantly in the singlet state (which is a bound state) and it will be incorrect to consider transition to unbound, plane wave states of  $c\bar{c}$  at this stage. Only at later stages, when the octet component is dominant, it may be appropriate to consider repulsive potential in Eqn.(12), and unbound  $c\bar{c}$  states for the transition probability. This means that the perturbation term should appropriately account for the growth of octet component for the potential in Eqn.(12), along with a continuing singlet component with corresponding singlet potential in Eqn.(12). This clearly is a complex issue, and a proper account of appropriate potential for this type of evolution of  $J/\psi$  cannot be carried out in simple approximation scheme considered here. We make a simplifying assumption that  $J/\psi$  becomes unbound only when it transforms to the octet representation *after* its interaction with the  $Z(3)$  wall. Until then it is assumed to be in the color singlet representation. Thus, in the calculations of the spatial excitation of the  $J/\psi$  state below, we use the  $c\bar{c}$  potential (Eqn.(12)) in the color singlet representation. The underlying physics is that incoming  $J/\psi$  is in the color singlet state, it interacts with  $Z(3)$  wall which excites it to higher state (spatial excitation), still

in color singlet potential. While traversing the  $Z(3)$  wall, and undergoing this spatial excitation, the  $J/\psi$  state also transforms to color octet state. The final state, after traversing the  $Z(3)$  wall, is spatially excited state in color octet representation, and our calculations give probability for this final state. This final octet state is unbound and hence such excited  $J/\psi$  disintegrates. We emphasize that at this stage, our aim is to point out the new possibilities of disintegration of  $J/\psi$  with  $Z(3)$  walls and this simplifying assumption should not affect our qualitative considerations and approximate estimates. We hope to give a more complete treatment in future. Thus, we continue to use the color singlet potential in Eqn.(12), while considering the spatial excitation of  $J/\psi$ .

Since the potential is central, we perform coordinate transformations

$$\vec{R}_{cm} = \frac{\vec{r}_1 + \vec{r}_2}{2} \quad \text{and} \quad \vec{r} = \vec{r}_1 - \vec{r}_2, \quad (13)$$

where,  $\vec{r}$  is the relative coordinate between  $q$  and  $\bar{q}$ .  $\vec{R}_{cm}$  is the center of mass of  $J/\psi$ . Using Eqn. (13) with Eqn. (1), we get

$$A_0^r = \gamma A_0^{11} [\gamma(z'_1 + vt')] - \gamma A_0^{11} [\gamma(-z'_2 + vt')]. \quad (14)$$

$z'_1$  and  $z'_2$  are written in terms of  $\vec{R}_{cm}$  and  $\vec{r}$ . Similar expressions can be obtained for  $A_0^b$  and  $A_0^g$ . In the above coordinates, the  $J/\psi$  wave function is  $\Psi(\vec{R}_{cm})\psi(\vec{r})$ . For simplicity, we assume that the center of mass motion remains unaffected by the external perturbation. Then  $\Psi(\vec{R}_{cm})$  has the plain wave solution, while  $\psi(\vec{r})$  can be written  $\psi(r, \theta, \phi) = \psi(r)Y_l^m(\cos \theta, \phi)$ . As  $J/\psi$  is the  $l = 0$  state, we have

$$\psi_i = \psi(r)Y_0^0 \quad \text{and} \quad \psi_j = \psi_n(r)Y_l^m(\cos \theta, \phi). \quad (15)$$

The radial part,  $\psi(r)$ , is obtained by solving radial part of Schrödinger equation with effective potential given by

$$V(r) = -\frac{\alpha_s C_F}{r} + C_{cnf} \sigma r + \frac{l(l+1)}{2\mu r^2} \quad (16)$$

where  $\mu$  is the reduced mass. When we use Eqn. (11), (14) and (15) in Eqn. (3), we get one of the terms as

$$\int_{-\infty}^{\infty} \psi_j^* A_0^r \psi_i d\vec{r}_1 d\vec{r}_2 = \int_0^{\infty} \int_{-1}^1 \int_0^{2\pi} \psi_n^*(r) Y_l^{m*}(\cos \theta, \phi) A_0^r Y_0^0 \psi_{100}(r) r^2 dr d(\cos \theta) d\phi. \quad (17)$$

In the above equation, we have ignored the motion of the center of mass of charmonium and have considered only the relative coordinate. Under  $\cos \theta \rightarrow -\cos \theta$ ,  $A_0^r \rightarrow -A_0^r$  and  $\psi_i$  does not change. So if  $Y_l^m(\cos \theta, \phi) = Y_l^m(-\cos \theta, \phi)$  then RHS of Eqn. (17) is zero. Thus we do not get any transition to a state which

is symmetric under  $\cos\theta \rightarrow -\cos\theta$ . This has very important significance. While the color part prohibits the transition to singlet final states, the space dependence of interaction forbids the transition to the  $l = 0$  state (in color octet). Thus we see that purely color excitation of  $J/\psi$  due to  $A_0$  field of a domain wall is not possible. The excitation is possible to the first excited state of an octet (like an ‘octet  $\chi$ ’ state). As the excited state will have a radius larger than the  $l = 0$  state it is more prone to melting in the medium, (though with color octet composition, the final state becomes unbound anyway).

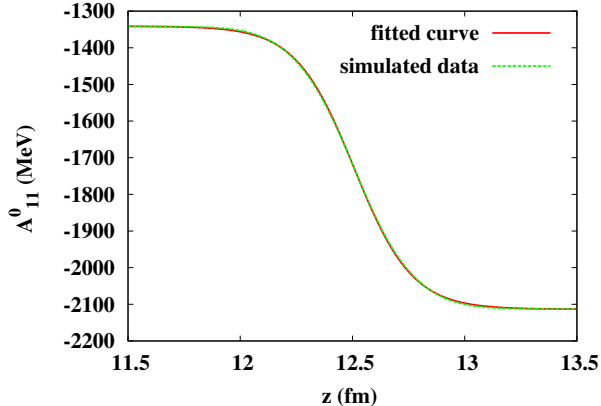


FIG. 1: (Color online)  $A_0$  profile across the  $Z(3)$  domain wall for  $T = 400$  MeV. Only  $(1,1)$  component is shown. Other components are similar. See ref. [19] for details.

### III. RESULTS

We numerically compute the integral given in Eqn. (3) with various parameters given after Eqn.(12). The profile of  $A_0$  is calculated from the profile of the Polyakov loop order parameter for a  $Z(3)$  domain wall at a temperature  $T = 400$  MeV (as a sample value). The details of this are given in ref. [19]. As explained there, the resulting profile is very well fitted by the functional form  $p \tanh(qx+r)+s$ , see Fig. 1.

We calculated the wave functions for various states of  $c\bar{c}$  with the complete potential given by Eqn. (16). For the calculation of the wave-functions for various states of  $c\bar{c}$  we have used Numerov method for solving the Schrödinger equation. We have also used energy minimization technique to get the wave functions and the bound state energy and the results obtained by both the methods match very well. Fig. (2) shows the radial part of the wave function for the  $l = 0, 1$  states of charmonium. The bound state contributions to the energy (excluding the rest mass of quarks) are found to be  $E_0 = 0.447$  GeV for  $J/\psi$  and  $E_0 = 0.803$  GeV for  $\chi$  state ( $l = 1$ ). We see from Fig.2 that the radius of  $J/\psi$  is about  $0.5$  fm while that for  $\chi$  is about  $0.8$  fm. Debye length in QGP

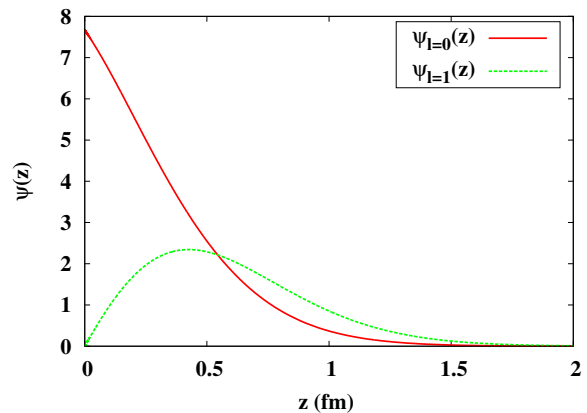


FIG. 2: (Color online) Wave functions for  $J/\psi$  ( $l = 0$ ) and  $\chi$  ( $l = 1$ ) states.

at  $T = 200$  MeV is  $r_d \sim 0.6$  fm and smaller at higher temperatures. Thus  $\chi$  state is unstable and it should melt easily in the medium (apart from the fact that in color octet state it also becomes unbound). Fig.3 shows the combined probability of transition to both the color octet  $\chi$  states (Eqns.(9),(10)) for an incoming  $J/\psi$  with different velocities moving normal to the domain wall. As we see, the probability rapidly rises as a function of velocity. However, for large velocities the probability of transition becomes large making first order perturbation approximation insufficient, and one needs more reliable estimates. Thus, the plot in Fig.3 should be trusted only for small velocities. Nonetheless, the trend at higher velocities strongly suggests that most of  $J/\psi$  will disintegrate while interacting with  $Z(3)$  walls.

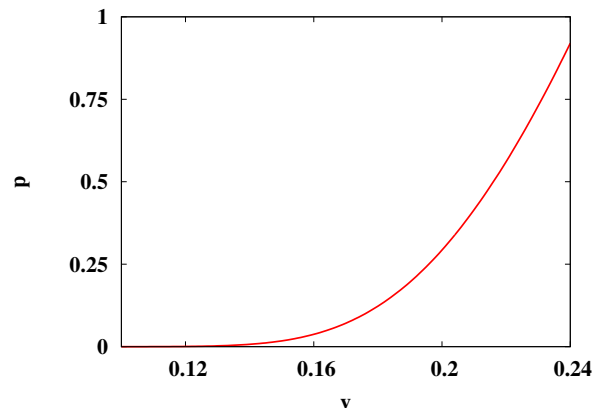


FIG. 3: (Color online) Probability  $p$  of transition of  $J/\psi$  to color octet  $\chi$  states vs. its velocity  $v$ . Note that the probability rapidly rises with  $v$ .

#### IV. CONCLUSIONS

These results show that on interaction with a  $Z(3)$  domain wall, a  $J/\psi$  particle will make an excitation to a higher orbital state in color octet representation which is unbound and will readily melt in the surrounding QGP medium. At higher energies, the transition probability keeps increasing, making the first order perturbation theory inapplicable and the results are not trustworthy. Nonetheless, this implies that at higher energies, almost all  $J/\psi$  are expected to disintegrate in this manner. This strong  $P_T$  dependence of  $J/\psi$  disintegration probability is a distinctive signature of our model wherein the probability of disintegration of  $J/\psi$  is enhanced with higher  $P_T$ . This can be used to distinguish this mechanism from the conventional Debye screening suppression. A very crucial point in the entire discussion is the Debye screening of the  $A_0$  profile of the domain wall itself as it carries color. At temperature 400 MeV, the domain wall has a thickness of  $\sim 1.5$  fm and the Debye radius for QGP is  $\sim 0.7$  fm. This means that Debye screening will be effective outside a sphere of diameter  $\sim 1.5$  fm. So we do not expect the domain wall to be significantly Debye screened. In the above discussion, we have completely ignored the effects of a thermal bath (QGP medium) on the potential (Eqn. 12) between  $c\bar{c}$  ([22, 27]). However as these effects make the potential between  $c\bar{c}$  weaker, the charmonium state *swells*. So it will be even easier for the interaction to break these bound states. These temperature effects will also be crucial for other heavier  $q\bar{q}$  states like bottomonium as they have large binding energies. Another important aspect which has been ignored for the sake of

simplicity, in the above calculations, is the question of the center of mass motion. This assumption is correct only in an average sense as the average force ( $\Delta V/\Delta z$ ) acting on  $c$  and  $\bar{c}$  vanishes. This averaging is done over the thickness  $\Delta z$ , which is the thickness of the domain wall itself. However as the instantaneous force ( $\partial V/\partial z$ ) is non-zero, there is a non-zero instantaneous acceleration of the center of mass. A more detailed analysis of the problem is required to incorporate all these details. One also needs to include the effects of dynamical quarks leading to explicit breaking of  $Z(3)$  symmetry. We mention that such a disintegration of  $J/\psi$  from a color electric field may not necessarily come from a background domain wall arising in QGP medium. In a thermal medium there are always statistical fluctuations. These gluonic fluctuations will have energy of order  $\sim T$ . Depending on the correlation length of the fluctuation, a  $J/\psi$  passing through it may disintegrate via the mechanism discussed above. It would be interesting to study the effect of these thermal gluonic fluctuations on the spectrum of mesons.

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